

Semileptonic $\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2595)$ and $\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2625)$ decays in the molecular picture of $\Lambda_c(2595)$ and $\Lambda_c(2625)$

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We evaluate the partial decay widths for the semileptonic $\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2595)$ and $\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2625)$ decays from the perspective that these two Λ_c^* resonances are dynamically generated from the DN and D^*N interaction with coupled channels. We find that the ratio of the rates obtained for these two reactions is compatible with present experimental data and is very sensitive to the D^*N coupling, which becomes essential to obtain agreement with experiment. Together with the results obtained for the $\Lambda_b \rightarrow \pi^- \Lambda_c^*$ reactions, it gives strong support to the molecular picture of the two Λ_c^* resonances and the important role of the D^*N component neglected in prior studies of the $\Lambda_c(2595)$ from the molecular perspective.

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I. INTRODUCTION

The interaction of mesons with baryons using chiral dynamics and unitarity in coupled channels, the chiral unitary approach [1–5] has brought light into the nature of some baryonic resonances. The prediction of two states for the $\Lambda(1405)$ [3, 6] has been one example of it, and is now supported by experiments as shown in Refs. [7, 8] (see also note in the PDG concerning this issue [9]). In the charm sector the interaction of DN and coupled channels has also been considered [10, 11] and, as a consequence, the $\Lambda_c(2595)$ resonance is generated dynamically, bearing many analogies to the $\Lambda(1405)$, one of which states couples strongly to $\bar{K}N$. While for some time only pseudoscalar-baryon channels were used, at some point it became clear that the mixture of pseudoscalar-baryon and vector-baryon should be relevant in some cases. A first step in this direction was given in Ref. [12], followed by Refs. [13, 14] in the light sector and by Refs. [15, 16] in the charm sector. Concerning the $\Lambda_c(2595)$, the explicit consideration of the DN and D^*N channels, using pion exchange to connect them, is done in Ref. [15]. In Refs. [17, 18], SU(8) symmetry was used, with a symmetry breaking mechanism that gives rise to the Weinberg-Tomozawa interaction in the SU(3) sector. Both in Refs. [15] and [18], it was found that the $\Lambda_c(2595)$ ($J^P = 1/2^-$) couples strongly to DN and D^*N in s -wave. The $\Lambda_c(2625)$ ($J^P = 3/2^-$) was also found dynamically generated, coupling strongly to D^*N in s -wave.

In a recent paper [19], the $\Lambda_b \rightarrow \pi^- \Lambda_c(2595)$ and $\Lambda_b \rightarrow \pi^- \Lambda_c(2625)$ decays were studied, and it was found that they were very sensitive to the DN and D^*N couplings and to their relative sign. The experimental ratio of the branching fractions for the two decays was well reproduced with the results obtained in Ref. [15]. It was found that the coupling of $\Lambda_c(2595)$ to D^*N was essential to obtain agreement with experiment, and if the relative sign of the couplings was reversed there was a cancella-

tion of the DN and D^*N components that makes the $\Lambda_b \rightarrow \pi^- \Lambda_c(2595)$ partial decay width extremely small, in sheer disagreement with experiment.

Support for the picture of Refs. [15, 18] should come from accumulation of experimental data which can be reproduced by the models. In this respect, in the present work we want to show one such reaction, the $\Lambda_b \rightarrow (\bar{\nu}_l l) \Lambda_c^*$, with $\Lambda_c^* \equiv \Lambda_c(2595), \Lambda_c(2625)$. We develop here the formalism that provides the width for these decays within the molecular picture of Ref. [15] and show that the ratio of the branching fractions for these two reactions is compatible with present experimental data. Theoretical calculations for $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c$, with Λ_c the ground state, have been done before using constituent quark models [20] and QCD lattice simulation [21]. For Λ_c^* , this is the first calculation.

II. FORMALISM

The picture for the $\Lambda_b \rightarrow (l \bar{\nu}_l) \Lambda_c(2595)$ or $\Lambda_b \rightarrow (l \bar{\nu}_l) \Lambda_c(2625)$ reactions is given in Fig. 1 at the quark level.

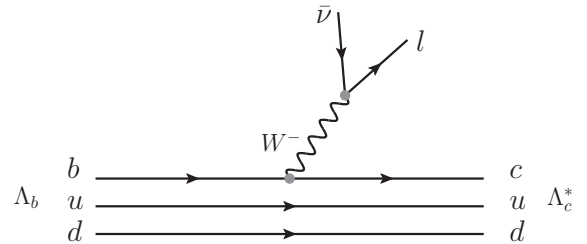


FIG. 1: Diagrammatic representation of the $\Lambda_b \rightarrow (l \bar{\nu}_l) \Lambda_c^*$ decay.

We must bear in mind three important points: 1) The ud quarks of the Λ_b are in $I = 0, S = 0$ and they are

spectators in the reaction. 2) Since in the final state, the ud quarks still have $I = 0$, $S = 0$ and positive parity, the c quark must carry negative parity to be able to produce the $1/2^-, 3/2^- \Lambda_c^*$ states at the end. This means it will have $L = 1$ in the quark picture. 3) Since the Λ_c^* is generated from the DN , D^*N interaction and other coupled channels, in the picture of Fig. 1 one must include hadronization, creating a $\bar{q}q$ pair with the quantum numbers of the vacuum. The coupling with $\bar{q}q$ to give a meson-baryon system must include the c quark to allow it to go back to the ground level, where it will be in the meson-baryon configuration.

The former considerations are similar to those done in the study of the $\Lambda_b \rightarrow J/\psi K^- p$ reaction studied in Ref. [22] and measured later in Ref. [23]. They were taken into account in the study of the $\Lambda_b \rightarrow \pi^- \Lambda_c^*$ decays in Ref. [19] and we make use of the results here. It was found there that after taking into account the hadronization, including a singlet of $SU(3)$ $\bar{q}q$ states ($\bar{u}u + \bar{d}d + \bar{s}s$), the following hadronic configuration appeared at the end, ignoring the larger mass $D_s^+ \Lambda$ component,

$$|H'\rangle = |D^0 p\rangle + |D^+ n\rangle \equiv \sqrt{2} |DN, I=0\rangle. \quad (1)$$

Similarly, the same combination of D^*N would appear, and the dynamics of the production of these two cases was explicitly studied in Ref. [19]. We shall use some of the findings of that work here.

The dynamics in the present case is different than the one found in the $\pi^- \Lambda^*$ decay. As shown in Refs. [24, 25], the transition matrix is given by

$$T = -iG_F \frac{V_{bc}}{\sqrt{2}} L^\alpha Q_\alpha V_{\text{had}}, \quad (2)$$

with G_F the Fermi coupling constant, V_{bc} the Cabibbo-Kobayashi-Maskawa matrix element for the $b \rightarrow c$ transition, V_{had} a factor accounting for the hadronic interaction, and L^α , Q_α the leptonic and quark operators,

$$L^\alpha \equiv \bar{u}_l \gamma^\alpha (1 - \gamma_5) v_\nu, \quad Q_\alpha \equiv \bar{u}_c \gamma_\alpha (1 - \gamma_5) u_b. \quad (3)$$

When evaluating the sum and average over polarizations of $|T|^2$, we will have

$$\frac{1}{2} \sum \sum |T|^2 \propto \frac{1}{2} \sum \sum |L^\alpha Q_\alpha|^2. \quad (4)$$

As shown in Ref. [24], we have

$$\begin{aligned} \sum_{\text{pol}} L^\alpha L^{\dagger\beta} &= \text{tr} \left[\gamma^\alpha (1 - \gamma_5) \frac{\not{p}_\nu - m_\nu}{2m_\nu} (1 + \gamma_5) \gamma^\beta \frac{\not{p}_l + m_l}{2m_l} \right] \\ &= 2 \frac{p_\nu^\alpha p_l^\beta + p_l^\alpha p_\nu^\beta - p_\nu \cdot p_l g^{\alpha\beta} - i\epsilon^{\rho\alpha\sigma\beta} p_{\nu\rho} p_{l\sigma}}{m_\nu m_l}. \end{aligned} \quad (5)$$

In Ref. [24], a sum and average over polarization of the quarks was also done for $Q_\alpha Q_\beta^\dagger$, but here we cannot do that if we want to differentiate between the production of

DN or D^*N . We divert from the formalism of Ref. [24] at this point, but recall from there that in the semileptonic processes the $\bar{u}l$ invariant mass is quite large, peaking around the end of the spectrum, which makes the Λ^* come out with relatively small momentum, and, sharing this momentum with the c, u, d quarks, the c quark carries a small momentum at the end compared to its mass. Using the nonrelativistic expressions for γ^μ and $\gamma^\mu \gamma_5$ and neglecting terms that go like p/m_c , we find that only the $\gamma^0 \simeq 1$, and the $\gamma^i \gamma_5 \sim \sigma^i (i = 1, 2, 3)$ components survive in this case. Then, after a bit of algebra, one easily finds

$$\begin{aligned} \sum_{\text{lepton pol.}} L^\alpha L^{\dagger\beta} Q_\alpha Q_\beta^\dagger &= \frac{2}{m_\nu m_l} [2p_\nu^0 p_l^0 - p_\nu \cdot p_l + 4p_\nu^0 \vec{p}_l \cdot \vec{\sigma} \\ &+ (\vec{p}_\nu \cdot \vec{\sigma})(\vec{p}_l \cdot \vec{\sigma}) + (\vec{p}_l \cdot \vec{\sigma})(\vec{p}_\nu \cdot \vec{\sigma}) + (p_\nu \cdot p_l)(\vec{\sigma} \cdot \vec{\sigma})], \end{aligned} \quad (6)$$

where $\vec{\sigma}$ is acting at the level of quarks and the proper matrix elements with the quark polarizations, and sum and average over them, are still to be done.

III. QUARK MATRIX ELEMENTS

The $\vec{\sigma}$ operators in Eq. (6) act on the spins of the b, c quarks, but, as mentioned above, the c quark is in $L = 1$. Then the quark matrix element that appears is

$$\begin{aligned} \mathcal{M} &\equiv \int d^3r \varphi_{\text{in}}(r) \varphi_{\text{out}}(r) e^{-i\vec{q} \cdot \vec{r}} \sum_m \mathcal{C}(1\frac{1}{2}J; m, M' - m) \\ &\times Y_{1m}^*(\hat{r}) \langle \frac{1}{2}, M' - m | O_P | \frac{1}{2}, M \rangle Y_{00}(\hat{r}), \end{aligned} \quad (7)$$

where Y_{1m}^* comes for the c quark and $Y_{00}(\hat{r})$ from the b quark, and $e^{-i\vec{q} \cdot \vec{r}}$ stands for the plane wave of the $\bar{u}l$ emitted pair with momentum \vec{q} . In Eq. (7), J is the total angular momentum of the c quark, which coincides with the spin of the Λ_c^* , $1/2$ for $\Lambda_c(2595)$ and $3/2$ for $\Lambda_c(2625)$, since the ud pair and the $\bar{q}q$ carry both $J' = 0$. The operator O_P will be 1 or $\vec{\sigma}$ depending on the terms in Eq. (6). Expanding $e^{-i\vec{q} \cdot \vec{r}}$ in partial waves, we have

$$e^{-i\vec{q} \cdot \vec{r}} = 4\pi \sum_l (-i)^l j_l(qr) \sum_\mu (-1)^\mu Y_{l\mu}^*(\hat{r}) Y_{l,-\mu}(\hat{q}). \quad (8)$$

After performing the $d\Omega(\hat{r})$ integration we get

$$\begin{aligned} \mathcal{M} &= -4\pi i \sum_m \mathcal{C}(1\frac{1}{2}J; m, M' - m) Y_{1m}^*(\hat{q}) \\ &\times \langle \frac{1}{2}, M' - m | O_P | \frac{1}{2}, M \rangle ME(q) Y_{00}, \end{aligned} \quad (9)$$

with

$$ME(q) \equiv \int r^2 dr \varphi_{\text{in}}(r) \varphi_{\text{out}}(r) j_1(qr). \quad (10)$$

In Ref. [19], the matrix elements \mathcal{M} were written in

terms of the macroscopic $\vec{\sigma}$ and \vec{S}^+ operators acting on the Λ_b and Λ_c^* , where \vec{S}^+ is the transition spin operator from spin $1/2$ to $3/2$, normalized such that

$$\sum_M S_i |M\rangle \langle M| S_j^+ = \frac{2}{3} \delta_{ij} - \frac{i}{3} \epsilon_{ijk} \sigma_k. \quad (11)$$

The results obtained are summarized in Table I, omitting the $ME(q)$ factor.

TABLE I: Macroscopic operators in the $\Lambda_b \rightarrow \Lambda_c^*$ transitions associated to the microscopic operators at quark level.

	$O_P = 1$	$O_P = \vec{\sigma} \cdot \vec{q}$ (quark level)
$J = 1/2$	$i \frac{\vec{\sigma} \cdot \vec{q}}{q}$	iq
$J = 3/2$	$-i\sqrt{3} \frac{\vec{S}^+ \cdot \vec{q}}{q}$	0

Coming back to Eq.(6), let us evaluate these terms now. We can take advantage that for the Λ_b the spin of the b quark is the same as the one of the Λ_b , since the ud pair comes $S = 0$. Then, in the sum over the third component of the Λ_b spin, M , we would have

$$\sum_M \sigma^i \left| \frac{1}{2} M \right\rangle \left\langle \frac{1}{2} M \right| \sigma^i \equiv \delta_{ii} + i\epsilon_{ijk} \sigma_k = 3, \quad (12)$$

and

$$\begin{aligned} & \sum_M \left[p_{\nu_i} \sigma_i \left| \frac{1}{2} M \right\rangle \left\langle \frac{1}{2} M \right| p_{l_j} \sigma_j \right. \\ & \quad \left. + p_{l_i} \sigma_i \left| \frac{1}{2} M \right\rangle \left\langle \frac{1}{2} M \right| p_{\nu_j} \sigma_j \right] \\ &= (p_{\nu_i} p_{l_j} + p_{\nu_j} p_{l_i}) (\delta_{ij} + i\epsilon_{ijk} \sigma_k) \\ &= 2\vec{p}_\nu \cdot \vec{p}_l. \end{aligned} \quad (13)$$

On the other hand, we can write the term $\vec{p}_l \cdot \vec{\sigma}$ as

$$\vec{p}_l \cdot \vec{\sigma} = \frac{1}{2} (\vec{q} + \vec{p}_r) \cdot \vec{\sigma}$$

where

$$\vec{q} = \vec{p}_l + \vec{p}_r, \quad \vec{p}_r = \vec{p}_l - \vec{p}_\nu.$$

Now, according to Table I, the term $p_\nu^0 \vec{\sigma} \cdot \vec{q}$ at the quark level, containing the operators 1 and $\vec{\sigma} \cdot \vec{q}$, will give rise to $p_\nu^0 i \frac{\vec{\sigma} \cdot \vec{q}}{q} (-)iq$ at the macroscopic level for $J = 1/2$ and zero for $J = 3/2$. But the trace of $\vec{\sigma} \cdot \vec{q}$, when summing over polarizations will be zero. The term $\vec{p}_r \cdot \vec{\sigma}$ will also vanish when one integrates over the angles, as we shall see in the next section. Hence, this term vanishes in the sum over polarizations and integration over the phase space. We are thus left with

$$\begin{aligned} & \sum_{\text{lepton pol.}} L^\alpha L^{\dagger\beta} Q_\alpha Q_\beta^\dagger \\ &= \frac{2}{m_\nu m_l} [2p_\nu^0 p_l^0 - p_\nu \cdot p_l + 2\vec{p}_\nu \cdot \vec{p}_l + 3p_\nu \cdot p_l] \\ &= \frac{8}{m_\nu m_l} p_\nu^0 p_l^0. \end{aligned} \quad (14)$$

Now if we want to evaluate the extra sum over polarizations of Λ_b and Λ_c^* , we go to the macroscopic representation of Table I and have

$$\overline{\sum} \sum L^\alpha L^{\dagger\beta} Q_\alpha Q_\beta^\dagger \equiv \frac{8}{m_\nu m_l} p_\nu^0 p_l^0 \frac{1}{2q^2} \sum_{M, M'} \begin{cases} \langle \frac{1}{2} M | \vec{\sigma} \cdot \vec{q} | \frac{1}{2} M' \rangle \langle \frac{1}{2} M' | \vec{\sigma} \cdot \vec{q} | \frac{1}{2} M \rangle, & \text{for } J = 1/2; \\ 3 \langle \frac{1}{2} M | \vec{S}^+ \cdot \vec{q} | \frac{3}{2} M' \rangle \langle \frac{3}{2} M' | \vec{S}^+ \cdot \vec{q} | \frac{1}{2} M \rangle, & \text{for } J = 3/2. \end{cases} \quad (15)$$

Hence,

$$\overline{\sum} \sum L^\alpha L^{\dagger\beta} Q_\alpha Q_\beta^\dagger \equiv A_J \frac{8}{m_\nu m_l} p_\nu^0 p_l^0, \quad (16)$$

with $A_{1/2} = 1$ and $A_{3/2} = 2$.

There is still one more element to consider, which is

to include the molecular dynamics of the Λ_c^* states. To connect to the DN and D^*N components one must take into account the hadronization of the $\bar{q}q$ pair. This was done in Ref. [19]. The mechanism is depicted in Fig. 2 and introduces a hadronic factor different for DN and D^*N coupling to $J = 1/2$, and for D^*N coupling to

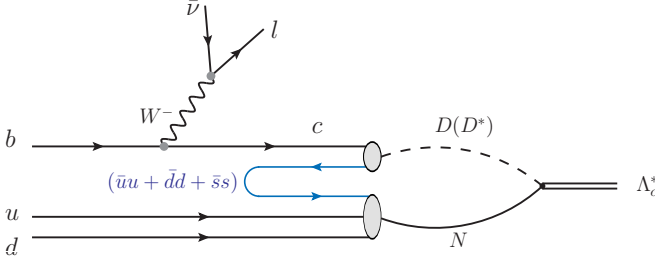


FIG. 2: Mechanism to produce a Λ_c^* resonance that is mostly made from DN, D^*N .

$J = 3/2$. The result of the hadronic factor V_{had} for these cases is given in Table II, up to a common global factor, where G_{DN} and G_{D^*N} are the $DN(D^*N)$ loop functions for the propagation of these states in Fig. 2, and $g_{R,DN}$ and g_{R,D^*N} the couplings of the $\Lambda_c(2595)$, $\Lambda_c(2625)$ to the DN and D^*N states. All this information is given in Ref. [15] and we summarize it in Table III.

TABLE II: Contributions to V_{had} from DN and D^*N in the coupling to $J = 1/2$ and $J = 3/2$ from Ref. [19]. The couplings g_i and the loop functions G_i are obtained in Ref. [15].^a

V_{had}	DN	D^*N
$J = 1/2$	$\frac{1}{2}G_{DN} \cdot g_{R,DN}$	$\frac{1}{2\sqrt{3}}G_{D^*N} \cdot g_{R,D^*N}$
$J = 3/2$	0	$\frac{1}{\sqrt{3}}G_{D^*N} \cdot g_{R,D^*N}$

^aNote change of sign in the D^*N case, as discussed in Ref. [19], because in Ref. [15] V_{eff}^2 for the $DN \rightarrow D^*N$ transition was calculated and the positive sign for V_{eff} was taken by default. The right sign, corresponding to π exchange, is negative. This sign is not relevant for the spectrum discussed in Ref. [15], but it matters here.

TABLE III: The values of $G_{DN} \cdot g_{R,DN}$ and $G_{D^*N} \cdot g_{R,D^*N}$ for the two Λ_c^* resonances.

	$G_{DN} \cdot g_{R,DN}$	$G_{D^*N} \cdot g_{R,D^*N}$
$\Lambda_c(2595)$	$13.88 - 1.06i$	$26.51 + 2.1i$
$\Lambda_c(2625)$	0	29.10

With all these ingredients, we can write

$$\sum \sum |T|^2 = C \frac{8}{m_\nu m_l} p_\nu^0 p_l^0 A_J V_{\text{had}}(J), \quad (17)$$

$$A_J V_{\text{had}}(J) \equiv \begin{cases} \left| \frac{1}{2}G_{DN} \cdot g_{R,DN} + \frac{1}{2\sqrt{3}}G_{D^*N} \cdot g_{R,D^*N} \right|^2, & \text{for } J = 1/2 \\ 2 \left| \frac{1}{\sqrt{3}}G_{D^*N} \cdot g_{R,D^*N} \right|^2, & \text{for } J = 3/2 \end{cases} \quad (18)$$

where C is a global common factor that contains $ME(q)^2$. With values of $M_{\text{inv}}(\bar{\nu}l)$ large, the values of q are not large. We can consider it a smooth function over the phase space. However, since the only observable that we want to evaluate is the ratio of branching fractions, this ratio is essentially given by the ratio of $A_J V_{\text{had}}(J)$ for the two resonances since the integrals of phase space are practically identical for the two resonances. Before we perform the numerical calculations of the phase space, we can already quote here that

$$\frac{A_{1/2} V_{\text{had}}(1/2)}{A_{3/2} V_{\text{had}}(3/2)} = 0.38, \quad (19)$$

and this should be very similar to the final result considering the slight differences in the phase space.

IV. EVALUATION OF THE WIDTH

The width for the decay into three particles is given by

$$\Gamma = 2M_{\Lambda_c^*} 2m_\nu 2m_l \frac{1}{2} \frac{1}{(2\pi)^3} \frac{1}{M_{\Lambda_b}} \times \int dM_{\text{inv}}(\bar{\nu}l) p_{\Lambda_c^*} d\Omega_{\Lambda_c^*} \int d\tilde{\Omega}_l \frac{1}{16\pi^2} \tilde{p}_l \sum \sum |T|^2, \quad (20)$$

where $p_{\Lambda_c^*} = -\vec{q}$ is the momentum of the Λ_c^* in the rest frame of Λ_b , and \tilde{p}_l the lepton momentum in the $\bar{\nu}l$ rest frame, respectively. One can see that the masses of the lepton and neutrino, which appear because of our choice of normalization of the fermion field, cancel in the final expression of Eq. (20). Once we arrive to this point we can come back to see why the $\vec{\sigma} \cdot \vec{p}_r$ term vanishes in the phase space integration. To see that, it is interesting to make a boost from the $\bar{\nu}l$ rest frame to the Λ_b rest frame where the $\bar{\nu}l$ pair has an energy $E_{\nu l}$ and a momentum \vec{q} . We obtain

$$\vec{p}_l = \left[\left(\frac{E_{\nu l}}{M_{\text{inv}}} - 1 \right) \frac{\vec{p}_l \cdot \vec{q}}{\vec{q}^2} + \frac{\vec{p}_l^0}{M_{\text{inv}}} \right] \vec{q} + \vec{p}_l, \quad (21)$$

where \vec{p}_l and $\vec{\tilde{p}}_l$ are the lepton momenta in the Λ_b rest frame and in the $\bar{\nu}l$ rest frame, respectively, and \vec{p}_l^0 the lepton energy in the $\bar{\nu}l$ rest frame. Since $\vec{p}_\nu = -\vec{p}_l$, then

$$\vec{p}_r = \vec{p}_l - \vec{p}_\nu = 2\vec{p}_l + 2 \left(\frac{E_{\nu l}}{M_{\text{inv}}} - 1 \right) \frac{\vec{p}_l \cdot \vec{q}}{\vec{q}^2} \vec{q}, \quad (22)$$

where we have considered that $\vec{E}_\nu = \vec{E}_l$, assuming zero mass for both of them.

We can see in Eq. (22) that when integrating over $d\tilde{\Omega}_l$ the vector \vec{p}_r , proportional to $\vec{\tilde{p}}_l$, will vanish in the integration.

One last point is a practical one to reduce the integral of Γ in Eq. (20) to just one numerical integration. For this we follow the steps of Ref. [24].

The factor $p_\nu^0 p_l^0$ in Eq. (16) evaluated in the Λ_b rest frame, where we could reduce the $\gamma^\mu, \gamma^\mu \gamma_5$ matrices to easy expressions, will depend on angles and we should in principle perform all the integrals in Eq. (20). We can write in an invariant way

$$p_\nu^0 p_l^0 (\Lambda_b \text{ rest frame}) = \frac{1}{M_{\Lambda_b}^2} (p_\nu \cdot p_{\Lambda_b}) (p_l \cdot p_{\Lambda_b}). \quad (23)$$

Next we evaluate these invariant products in the frame where $\bar{\nu}l$ is at rest. In this frame, $\vec{p}_l = -\vec{p}_\nu$, $\vec{p}_{\Lambda_b} = \vec{p}_{\Lambda_c^*}$ and using $\vec{E}_{\Lambda_b} = M_{\text{inv}} + \vec{E}_{\Lambda_c^*}$ we obtain

$$\vec{E}_{\Lambda_b} = \frac{M_{\Lambda_b}^2 + M_{\text{inv}}^2 - M_{\Lambda_c^*}^2}{2M_{\text{inv}}}, \quad |\vec{p}_{\Lambda_b}| = \frac{\lambda^{1/2}(M_{\Lambda_b}^2, M_{\text{inv}}^2, M_{\Lambda_c^*}^2)}{2M_{\text{inv}}}.$$

Then,

$$\begin{aligned} & \frac{1}{M_{\Lambda_b}^2} (p_\nu \cdot p_{\Lambda_b}) (p_l \cdot p_{\Lambda_b}) \\ &= \frac{1}{M_{\Lambda_b}^2} \left[(\vec{p}_\nu \cdot \vec{E}_{\Lambda_b} - \vec{p}_\nu \cdot \vec{p}_{\Lambda_b}) (\vec{p}_l \cdot \vec{E}_{\Lambda_b} - \vec{p}_l \cdot \vec{p}_{\Lambda_b}) \right] \\ &= \frac{1}{M_{\Lambda_b}^2} \left[(\vec{p}_\nu \cdot \vec{E}_{\Lambda_b})^2 - (\vec{p}_\nu \cdot \vec{p}_{\Lambda_b})^2 \right]. \end{aligned} \quad (24)$$

Taking into account that

$$\int d\tilde{\Omega}_l |\vec{p}_l| |\vec{p}_{\Lambda_b}|^2 \cos^2 \theta = \frac{1}{3} \int d\tilde{\Omega}_l |\vec{p}_l| |\vec{p}_{\Lambda_b}|^2,$$

and that

$$\tilde{p}_\nu^0 = \frac{M_{\text{inv}}}{2} = |\vec{p}_l|,$$

we obtain, that we can neglect the angle dependence of $p_\nu^0 p_l^0$ and use over the whole phase space

$$p_\nu^0 p_l^0 \rightarrow \frac{1}{M_{\Lambda_b}^2} \left(\frac{M_{\text{inv}}}{2} \right)^2 \left[\tilde{E}_{\Lambda_b}^2 - \frac{1}{3} \tilde{p}_{\Lambda_b}^2 \right], \quad (25)$$

and the width is now given by

$$\Gamma = \int \frac{d\Gamma}{dM_{\text{inv}}} dM_{\text{inv}} \quad (26)$$

with

$$\frac{d\Gamma}{dM_{\text{inv}}} = 2M_{\Lambda_b} 2M_{\Lambda_c^*} 2m_\nu 2m_l \frac{1}{4M_{\Lambda_b}^2} \frac{1}{(2\pi)^3} p_{\Lambda_c^*} \tilde{p}_l \sum \sum |T|^2, \quad (27)$$

where in $|T|^2$ of Eq. (17) we substitute $p_\nu^0 p_l^0$ by the expression of Eq. (25), and $p_{\Lambda_c^*}, \tilde{p}_l$ are given by

$$\begin{aligned} p_{\Lambda_c^*} &= \frac{\lambda^{1/2}(M_{\Lambda_b}^2, M_{\text{inv}}^2, M_{\Lambda_c^*}^2)}{2M_{\Lambda_b}}, \\ |\tilde{p}_l| &= \frac{\lambda^{1/2}(M_{\text{inv}}^2, m_l^2, m_\nu^2)}{2M_{\text{inv}}} \equiv \frac{M_{\text{inv}}}{2}. \end{aligned}$$

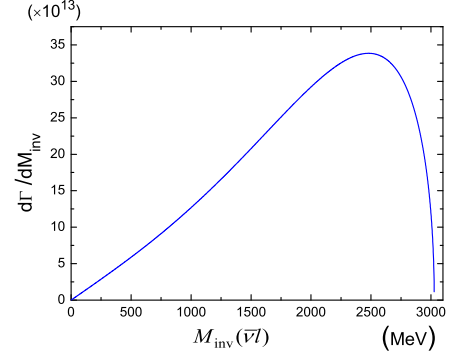


FIG. 3: $\frac{d\Gamma}{dM_{\text{inv}}}$ for the $(\bar{\nu}l)$ pair as a function of $M_{\text{inv}}(\bar{\nu}l)$ in the $\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2595)$ decay.

V. RESULTS

We evaluate first $\frac{d\Gamma}{dM_{\text{inv}}}$ for the $\Lambda_c(2595)$ by means of Eq. (27), and the results are shown in Fig. 3. The result for $\Lambda_c(2625)$ has a nearly identical shape. As we can see from Fig. 3, the $(\bar{\nu}l)$ invariant mass distribution peaks around the end of the phase space. Thanks to this, the momenta of Λ_c^* are relatively small, justifying the non-relativistic approximations made in γ^μ and $\gamma^\mu \gamma_5$ in the Wbc vertex.

On the other hand, by taking C constant in Eq. (17) and using Eqs. (17), (25), (27), we evaluate $\Gamma[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2595)]$ and $\Gamma[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2625)]$. We eliminate C by taking the ratio of the two widths, and we find

$$\frac{\Gamma[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2595)]}{\Gamma[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2625)]} = 0.39. \quad (28)$$

As we can see, this result is practically identical to the one obtained in Eq. (19). The effect of considering the different phase space in the two reactions of Eq. (28) is an increase of the ratio by 3% with respect to the result obtained in Eq. (19).

The experimental data from the PDG are [26]

$$\begin{aligned} BR[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2595)] &= (7.9_{-3.5}^{+4.0}) \times 10^{-3}, \\ BR[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2625)] &= (13.0_{-5.0}^{+6.0}) \times 10^{-3}. \end{aligned} \quad (29)$$

The ratio, summing in quadrature the experimental errors is

$$\frac{\Gamma[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2595)]}{\Gamma[\Lambda_b \rightarrow \bar{\nu}_l \Lambda_c(2625)]} \Big|_{\text{Exp.}} = 0.6_{-0.3}^{+0.4}. \quad (30)$$

We can see that there is agreement between theory and experiment within errors.

The agreement obtained is not trivial and essentially tied to the D^*N component of the $\Lambda_c^*(2595)$ resonance. Should there be no coupling to D^*N , we would have obtained a ratio for Eq. (28) of the order of 0.1, clearly in contradiction with experiment, even within the large er-

rors. On the other hand, should the relative sign between g_{R,D^*N} and $g_{R,DN}$ be the opposite, we would have obtained a ratio for Eq. (28) of 0.02 in sheer contradiction with experiment.

The reactions studied and their ratio of widths give support to the molecular picture of the $\Lambda_c(2595)$ and $\Lambda_c(2625)$ as dynamically generated from DN, D^*N and other coupled channels, described in Ref. [15], with DN and D^*N as the more important components. Together with the results obtained in Ref. [19] for the $\Lambda_b \rightarrow \pi^- \Lambda_c(2595)$ and $\Lambda_b \rightarrow \pi^- \Lambda_c(2625)$, they provide a boost to this molecular picture. It would be good to have evaluations of these ratios with different pictures, as well as have more experiments with the production of these resonances which can be contrasted with the different pictures.

VI. CONCLUSIONS

We have studied the $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2595)$ and $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c(2625)$ reactions from the perspective that the two Λ_c^* resonances are dynamically generated from the DN, D^*N interaction with coupled channels. We work out microscopically the weak vertices, involving Wbc and $W\bar{\nu}l$, to have a DN, D^*N baryonic final state, which couples to Λ_c^* . For this, a $\bar{q}q$ pair with the quantum numbers of the vacuum is created and the $c\bar{q}$ combine to give either the D or D^* . With the help of Racah algebra, one can work out the weight for the formation of the DN and D^*N components. This, together with the coupling of DN, D^*N to the two Λ_c^* states, gives finally the amplitudes for the $\Lambda_b \rightarrow \bar{\nu}_l l \Lambda_c^*$ transitions. With the input for the DN, D^*N couplings to Λ_c^* obtained in Ref. [15] we can evaluate the rates for these transitions, up to a common factor involving radial matrix elements of the b and c

wave functions. We do not evaluate this matrix element, which involves an excited c quark state in $L = 1$, but calculate the ratio of partial decay widths, where this factor cancels. We obtain results which are in agreement with experiment, within errors, and note that the agreement is obtained thanks to the coupling of the $\Lambda_c(2595)$ to the D^*N component, which was neglected in early studies of this resonance. This agreement adds to the one found before for the $\Lambda_b \rightarrow \pi^- \Lambda_c^*$ reactions. One should note that the ratio found for these latter branching fractions ($\Lambda_c(2595)$ versus $\Lambda_c(2625)$) was about 0.74, while the one for the semileptonic reactions has been found of the order of 0.4. The experimental data also follow this trend, the ratio for $\pi^- \Lambda_c^*$ decays is of the order of 1.0 ± 0.6 , while the one of the semileptonic decays is about 0.6 ± 0.4 . The relative weight of these ratios for the central values is very similar in the theory and the experiment.

The results obtained with these reactions give support to the molecular picture for these two Λ_c^* resonances. Work with other models and checks for further experiments will help us gain further insight on the nature of these resonances, and new experiments producing these two resonances should be encouraged.

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